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Strange Asymmetries in the Nucleon Sea

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Abstract

Relativistic corrections to the strange axial form factor are evaluated within a light-cone formalism, taking into account effects due to the breaking of Lorentz covariance associated with the use of one-body currents. Similar corrections are known to also be needed for the magnetic form factor, and we discuss the importance of these for recent data on strange electromagnetic form factors from the HAPPEX Collaboration at Jefferson Lab. The strange vector form factors, the strange axial charge, and the asymmetries in strange quark distributions are shown to be consistently correlated within the chiral cloud model.

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I. INTRODUCTION

The sea of the nucleon continues to intrigue nuclear and particle physicists seeking to understand its structure and dynamical origin. The recent results from the E866 Collaboration [1] at Fermilab, for example, on Drell-Yan production in proton–proton and proton–deuteron scattering, from which the x-dependence of the $\bar{d}/\bar{u}$ ratio was extracted, present a challenge to models of the nucleon's structure [2]. These data quite unambiguously indicate non-trivial non-perturbative effects in the proton's sea which give rise to a rather large asymmetry in the light antiquark sector.

At somewhat lower energies, the HAPPEX Collaboration [3] has recently presented results on the strange electromagnetic form factors of the proton obtained from parity-violating electron scattering at Jefferson Lab. The experiment found $G_E^S + rG_M^S\Big|_{\text{(HAPPEX)}} \approx 0.023 \pm 0.048$ at an average Q^2 of 0.48 GeV^2, with $r \sim 0.4$ for the HAPPEX kinematics. This result is consistent with the earlier experiment by the SAMPLE Collaboration at MIT-Bates [4], $G^S_{M\text{(SAMPLE)}} = +0.23 \pm 0.44$ at $Q^2 = 0.1$ GeV^2.

These experiments are extremely valuable in our quest to arrive at a consistent picture of the nucleon's substructure. While valence quark models have provided considerable insight into the structure of the nucleon's core, describing the dynamics of the sea of the nucleon is considerably more model-dependent. Nevertheless, the nucleon sea provides a unique testing ground for QCD models, since a sea generated purely perturbatively generally results in vanishing sea quark form factors and asymmetries.

While constrained by conservation laws requiring equal numbers of strange and anti-strange quarks in the nucleon (which in deep-inelastic scattering language corresponds to equal first moments of the s and $\bar{s}$ quark distributions, or in elastic scattering, a zero strange electric form factor at $Q^2 = 0$), the distributions of s and $\bar{s}$ quarks need not be identical in coordinate or momentum space [5,6]. Indeed, since perturbative QCD predicts equal s and $\bar{s}$ distributions, a difference between these would be clear evidence for non-perturbative effects in the structure of the nucleon.

Many models have been constructed in the literature which attempt to describe how strangeness arises in the nucleon. These range from vector meson dominance and quark models, to Skyrme and NJL types, as well as approaches which try to respect general properties such as analyticity and chiral symmetry [7–15]. It is probably fair to say that none of these models is sophisticated enough or has the sufficient degrees of freedom necessary to provide a reliable microscopic description of all the strangeness observables. Nevertheless, in many cases such model studies can offer a glimpse of the underlying dynamics of strangeness generation.

A further complication arises with the need to consistently keep the same model degrees of freedom at different scales. For example, in deep-inelastic scattering, the natural degrees of freedom are partons on the light-cone; at low energies one can obtain reasonable descriptions of observables in terms of effective, constituent quarks. Until a rigorous connection is found between these (see however Ref. [16]), use of quark-type models will be problematic if one aims for a unified description of strangeness observables.

A somewhat less ambitious endeavour than calculating structure functions from first principles is to accept the limitations of the QCD models, and try to see whether a piece of strangeness information from one experiment can be used to understand data from another

experiment. This can also pose some challenges, as the validity of models is often limited to a specific energy range (for example, below the chiral symmetry breaking scale, for chiral hadronic models), forcing one to sometimes extrapolate models to regions where their reliability could be questioned.

One model which in the past has been applied to the study of both low energy observables, such as electromagnetic form factors, magnetic moments, hadron–hadron scattering, etc, as well as to deep-inelastic structure functions at much higher energies, is the chiral cloud model [17–20]. Since there is *a priori* no resolution scale at which chiral symmetry can be ignored, a cloud of pseudoscalar mesons would be expected to play some role both at low and high energies, provided one can isolate the non-perturbative effects from the purely perturbative ones associated with QCD evolution between different scales.

One of the main difficulties with the implementation of cloud models of the nucleon in the past has been how to evaluate matrix elements of current operators between non-physical states, such as the virtual mesons or baryons of the cloud. Some of these can be circumvented by formulating the cloud on the light-cone (or in the infinite momentum frame). The light-cone offers many advantages for the description of hadron and nuclear structure and interactions, as advocated some 20 years ago by Lepage & Brodsky [21] and others. Interpreting intermediate-state particles as being on their mass shells, one can avoid introducing *ad hoc* off-shell nucleon form factors or structure functions, and more consistently parameterise the momentum dependence of vertex functions at hadronic vertices. Furthermore, it is the natural framework for describing partonic substructure of hadrons and nuclei (see recent work by Miller and Machleidt [22] in the application of light-cone techniques to nuclear matter). Nevertheless, at low energies there are some subtle issues, such as rotational invariance, which need special care when dealing with models on the light-cone, and particular care is paid to these in this paper.

A common assumption in the application of chiral cloud and other models is the impulse approximation, in which one truncates the Fock space at some convenient point (usually determined by one's ability to calculate), and omits contributions from many-body currents. It is known, however, that the use of one-body currents alone for composite systems leads to a violation of Lorentz covariance [23,24]. In Section II we discuss the consequences of this, and in particular outline how Lorentz covariance may be restored using the prescription of Karmanov et al. [24–26]. More complete accounts can be found, for example, in Refs. [27,28]. In Section III the light-cone chiral cloud model is applied to the strange axial charge, and important corrections are found to arise from the Lorentz symmetry breaking effects. A similar analysis for the strange magnetic form factor was reported in Ref. [29], where the corrections changed even its sign, bringing it more into line with the SAMPLE measurement [4]. The results for the strange electromagnetic form factors from the chiral cloud model [29,30] are compared with the new data from the HAPPEX Collaboration [3] at Jefferson Lab in Section IV. Finally, in Section V we summarise our findings and outline possible improvements of our analysis in future work.

II. RELATIVISTIC COVARIANCE AND THE LIGHT-CONE

As is well known, the light-cone formulation of dynamics has many advantages over other formulations when dealing with composite systems. The more pertinent ones to the

current discussion are those connected with the fact that negative energy contributions to intermediate states are absent, and particles can be treated as if they were on-mass-shell. In practice, this means that matrix elements of hadrons (and nuclei) can be simply expressed as convolutions of the constituent particles' matrix elements, and the constituents' distributions in the hadron. The issue of off-mass-shellness has plagued many earlier, instant form calculations which attempted to incorporate relativistic effects [31]. This is not to say that the light-cone formulation solves all problems which arise in instant form approaches — rather, to some extent one merely reshuffles them according to what is most convenient for a given application. Furthermore, if one is to unambiguously correlate information on strangeness observables from different experiments, one must utilise the same framework consistently throughout. Since the light-cone is the appropriate framework for high-energy deep-inelastic scattering, it would appear natural to also use the same model on the light-cone to describing observables such as elastic form factors.

One should point out, however, that the issue of relativistic covariance is relevant both in light-cone [24] as well as instant-form [28] approaches, whenever a Fock state suffers some kind of truncation, which invariably leads to a violation of Lorentz covariance. The problem exists because one-body currents, which do not include interactions, and to which most model calculations are restricted, do not commute with the interaction-dependent generators of the Poincaré group. Consequently, an incorrect four-vector structure appears in the matrix elements of current operators, resulting in the presence of additional unphysical, or spurious, terms in a form factor expansion of any electroweak current matrix element. The spurious form factors would not be present if the Lorentz symmetry were exact.

In the explicitly covariant formulation of light-cone dynamics developed by Karmanov et al. [24,25], a specific method was proposed for extracting the nucleon's physical form factors, excluding the spurious contributions. In this formulation, the nucleon state vector is defined on a light-cone given by the invariant equation $n \cdot x = 0$, where n is an arbitrary light-like four vector, $n^2 = 0$. Since this formulation is covariant, the Lorentz symmetry is restored, but the matrix elements now depend on the position of the light-cone plane, n^μ, which in principle no physical quantity should. Because of the explicit dependence on the light-cone orientation, a form factor expansion of the matrix elements of the electroweak current will now involve three variables, the nucleon (p^μ) and photon (q^μ) four-momenta and n^μ, rather than the usual first two. Therefore in general more structures will appear in the form factor decomposition, some of whose coefficients (namely, those depending on n^μ) will be unphysical. The advantage of this prescription is that these n^μ dependent coefficients can then be identified and subtracted from the physical form factors.

One can also compare the covariant light-cone formulation with the approach more commonly used in the literature for calculating light-cone matrix elements, namely using the "+" component of currents, with $n^\mu = (1; 0, 0, -1)$, so that $t + z = 0$ defines the usual light-cone plane.

In Ref. [24] the corrections to the electromagnetic vector form factors of the nucleon were calculated in a quark model. It was found that while the electric form factor does not suffer from any contamination from spurious form factors, the magnetic form factor receives quite large contributions. Following a similar philosophy, the corrections to the strange vector form factors of the nucleon were estimated in Ref. [29] within a light-cone chiral cloud model. For intrinsically small quantities such as those involving strangeness,

any corrections are likely to be relatively more important, and indeed the strange magnetic form factor was seen to change sign when the spurious contributions were subtracted. In addition to the strange vector matrix elements, those of the strange axial vector current are also of considerable interest, as they convey information on the spin distribution of strange quarks in the nucleon, which has been actively debated since the discovery of the EMC-spin effect a decade ago [32]. In the following Section we examine the strange axial charge in the light-cone formulation, and estimate the contamination to this from the spurious form factors within the chiral cloud model.

III. STRANGE AXIAL CHARGE

Given that one of the main reasons for the focus on the strangeness content of the nucleon was the distribution of the proton's spin amongst its constituents, it is clearly important to test whether previous estimates of strange contributions to the axial charge are reliable. In this Section we apply the prescription described above to the nucleon's strange axial charge.

The strange axial current on the light-cone can be written covariantly as [25]:

$$J^S_{\mu 5} = g^S_A\ \gamma_\mu \gamma_5 + \ b^S_1\ \frac{\not{n} p_\mu}{n\cdot p}\gamma_5 + \ \cdots\ , \qquad (1)$$

where the "$\cdots$" represent terms which do not contribute to axial matrix elements. The b^S_1 in this decomposition arises precisely because of the extra n dependence introduced by the light-cone orientation. In an exact calculation it would be identically zero. In practice, however, when one uses Lorentz covariance violating approximations, such as restrictions to one-body currents, this n dependent form factor can be non-zero.

Taking the forward matrix element of the current $J^S_{\mu 5}$, one can extract the axial charge g^S_A by using the trace projection [25]:

$$g^S_A = \frac{1}{4(n\cdot p)^2}\mathrm{Tr}\left[\mathcal{O}_\mu\left((n\cdot p)p^\mu\ \not{n}\gamma_5 - M^2 n^\mu\ \not{n}\gamma_5 - (n\cdot p)^2\ \gamma^\mu\gamma_5\right)\right], \qquad (2)$$

where $\mathcal{O}_\mu = (\not{p}+M)J^S_{\mu 5}(\not{p}+M)/(4M^2)$. Without correcting for the unphysical b^S_1 form factor, the axial charge would be [25]:

$$\tilde{g}^S_A = g^S_A + b^S_1\ \ =\ \ \frac{M^2}{2(n\cdot p)^2}\mathrm{Tr}\left[\mathcal{O}_\mu\ n^\mu\ \not{n}\gamma_5\right]. \qquad (3)$$

To ascertain the importance of the difference between $\tilde{g}^S_A$ and the corrected g^S_A, one can use a simple chiral cloud model, in which the strangeness in the nucleon is assumed to reside in the kaon and hyperon components of the nucleon wave function. Because of the very different masses and momentum distributions of the kaon and hyperon, the overall strange and antistrange distributions will be quite different [5]. In particular, in the valence approximation for the cloud, the $\bar{s}$ distribution is expected to be zero, since it resides entirely in the scalar kaon.

In the chiral cloud model the nucleon couples to a pseudoscalar kaon (K) and a spin-1/2 hyperon (Y) via a pseudoscalar $i\gamma_5$ interaction (the same results are also obtained with a pseudovector coupling). Extension of this analysis to spin-3/2 hyperons or strange vector

mesons is straightforward, although beyond the scope of the present discussion (see below). Because the kaon has spin 0, the axial form factor receives contributions only from the $\gamma^*\Lambda$ coupling, which can be written:

$$g_A^S = \frac{g_{KNY}^2}{16\pi^3}\int \frac{dy\ d^2\mathbf{k}_T}{y^2(1-y)}\frac{\mathcal{F}^2}{(\mathcal{M}^2-M^2)^2}\left(1-\frac{M_Y}{yM}\right)\Big(k_T^2+M_Y(M_Y-yM)\Big)\,, \tag{4}$$

where $\mathcal{M}^2 = (k_T^2+M_Y^2)/y + (k_T^2+m_K^2)/(1-y)$ is the invariant mass of the intermediate state, and $\mathcal{F}$ parameterises the hyperon-meson-nucleon vertex.

The momentum dependence of the vertex function $\mathcal{F}$ can be calculated within the same model by dressing and renormalising the bare KNY vertex by K loops. However, since a detailed model description of the hadronic vertex is not the purpose of this paper, we shall instead follow the more phenomenological approach and parameterise the KNY vertex by a simple function, such as a monopole, $\mathcal{F} = (\Lambda^2+M^2)/(\Lambda^2+\mathcal{M}^2)$. We shall comment on the dependence of the strangeness distribution on the shape of the form factor later.

For the uncorrected strange axial charge, from Eq.(3) we have:

$$\widetilde{g}_A^S = \frac{g_{KNY}^2}{16\pi^3}\int \frac{dy\ d^2\mathbf{k}_T}{y^2(1-y)}\frac{\mathcal{F}^2}{(\mathcal{M}^2-M^2)^2}\Big(-k_T^2+(M_Y-yM)^2\Big)\,, \tag{5}$$

which agrees with the expressions obtained in Ref. [30]. The results for the strange axial charge g_A^S are shown in Fig.1 as a function of the cut-off mass Λ. In practice, the $K\Lambda$ configuration turns out to give by far the dominant contribution to g_A^S if standard coupling constants [33] and form factor cut-offs are used. Also shown in Fig.1 is the uncorrected charge $\widetilde{g}_A^S$. The n dependent form factor turns out to be rather large, and contaminates the "true" g_A^S to such an extent so as to produce the rather small $\widetilde{g}_A^S$ value observed. The only empirical information available on the strange axial charge comes from the Brookhaven 734 experiment [34] on elastic νp and $\bar{\nu} p$ scattering [35]. Unfortunately, the value of g_A^S extracted from this experiment was found to be strongly correlated with the value of the cut-off mass, M_A, in the dipole axial vector form factor parameterisation. Varying M_A between 1.086 ± 0.015 GeV and 1.012 ± 0.032 GeV, one can obtain anything between $g_A^S = 0$ and -0.21 ± 0.10 [36], as indicated by the shaded region in Fig.1.

One can also compare the strange axial charge with the first moment, Δs, of the polarised strange quark distribution measured in deep-inelastic scattering [37]. A recent world-averaged value extracted in the $\overline{\text{MS}}$ scheme at a scale of 10 GeV^2 is $\Delta s = -0.10 \pm 0.04$ [37], as indicated by the two long-dashed horizontal lines in Fig.1. Note that in the chiral cloud model the distribution Δs is given by a convolution of the y-integrand in Eqs.(4) and (5) with the polarised strange distribution in the hyperon [30]. Since the latter is not a δ-function, but has a non-trivial x-dependence, the resulting convolution would be expected to be smaller than the strange axial charge in the model. The experimental value of Δs is nonetheless consistent with the calculated g_A^S if a soft KNY form factor is used.

To constrain the size of the KNY form factor, which is essentially the only parameter in the chiral cloud model, one can compare the model predictions with the measured unpolarised s–$\bar{s}$ asymmetry. The possible differences between the s and $\bar{s}$ quark distributions in the nucleon were investigated by the CCFR Collaboration via charm production in ν and $\bar{\nu}$ deep-inelastic scattering [38]. Such differences were first predicted in the meson cloud framework more than 10 years ago by Signal and Thomas [5]. The x-dependence of

the calculated s–$\bar{s}$ distribution is shown in Fig.2 for a form factor cut-off mass of $\Lambda = 1$ GeV (which gives an average multiplicity of kaons in the nucleon of $\approx 6\%$). The shaded region represents the data from Ref. [38]. Also shown for comparison is the result with a t-dependent monopole form factor, as used in earlier analyses, $\mathcal{F} = (\Lambda - m_K^2)/(\Lambda - t)$, where $t = (p_N - p_Y)^2$. Notice that the final shape and sign of s–$\bar{s}$ are quite sensitive to the shape of the KNY form factor. On the other hand, it is known that form factors which depend solely on the t variable violate momentum conservation when one considers scattering from both the meson and hyperon components of the nucleon [39]. More precise measurement of the strange asymmetry would be a valuable test of the dynamics of the KNY interaction.

One can also compare the predictions of the model with the absolute values of the extracted s and $\bar{s}$ distributions, as done in Ref. [6] for example. As in Fig.2, one finds that for a hard hadronic form factor the meson cloud contributions overestimate the data, especially at large x [6]. The problem with comparing to the total s and $\bar{s}$ distributions, however, as distinct from their difference, is that the total distributions contain singlet contributions in addition to the non-singlet. Modeling the former in general requires the (symmetric) perturbative sea arising from $g \to s\bar{s}$, as well as additional input for the structure of the bare nucleon distributions, uncertainties in which consequently make any real predictions of the model more elusive. For this reason a comparison with the non-singlet difference $s - \bar{s}$, in which the perturbative contributions cancel, is more meaningful for the meson cloud model.

Finally, before ending this Section, we should note several concerns which have been raised in the literature regarding the implementation of loops in chiral models of the nucleon. In particular, it has been pointed out that truncations of the Fock state which stop at the one-loop order violate, in addition to the Lorentz covariance discussed above, also unitarity [15]. While this is true in principle, the region where rescattering should become an issue is above the production threshold, which in practice is at rather high momenta compared with those most relevant to the current process. Furthermore, the chiral cloud model discussed here rests on a perturbative treatment of the effective hadronic Lagrangian, so that provided the form factors used at the hadronic vertices are not very hard, one would expect a one-loop calculation for the most part to give the dominant contribution. If two loop contributions were found to be large compared with the leading ones, the perturbative formulation of the chiral cloud itself would need to be reconsidered. Recent work [40] based on coherent states techniques which include effects of higher-order, multi-pion Fock states for models such as the cloudy bag [17] indicates that for relatively small meson densities, a one-loop, perturbative treatment comes very close to the exact result. The conclusion is therefore that so long as the hadronic vertices are relatively soft, with $\Lambda \sim 1$ GeV, the one-loop result should give a reasonable estimate of cloud effects.

Concerns have also been raised about the omission of contributions from higher-mass intermediate states in the meson–baryon fluctuations [41]. While the effects of heavier baryons such as the Σ^* have been shown to be negligible [30], it has been argued that strange vector meson contributions are of the same order of magnitude as the K. In the analysis of Ref. [41] a rather hard $K^*N\Lambda$ form factor was used, however, with a cut-off mass in the monopole parametrisation of $\Lambda \sim 2.2$ GeV. This is to be compared with $\Lambda \sim 1.2$ GeV for the $KN\Lambda$ vertex [41]. This relatively large value for the $K^*N\Lambda$ cut-off was taken from the hyperon–nucleon scattering analysis of Ref. [33], although more recent work [42,43] suggests that a value for *both* the $K^*N\Lambda$ and $KN\Lambda$ form factor cut-offs of ~ 1 GeV is

more appropriate, Such a smaller value would significantly reduce the K^* contribution. A re-evaluation of the strange vector meson effects with softer form factors would therefore be very useful before definitive conclusions about the reliability of lowest-order one-loop calculations can be made.

IV. STRANGE ELECTROMAGNETIC FORM FACTORS

As well as understanding polarised strangeness in the nucleon, there has also been considerable effort directed at measuring matrix elements of the electromagnetic vector currents. The first experimental result on the strange magnetic form factor of the proton was obtained by the SAMPLE Collaboration [4] at MIT-Bates in 1997 in parity-violating electron scattering at backward angles, at $Q^2 = 0.1$ GeV^2. While plagued with large errors, the data did seem to favour a relatively small, and possibly positive, value of the strange magnetic moment. More recently, the HAPPEX Collaboration at Jefferson Lab [3] performed a similar experiment, although at forward angles, measuring the left-right asymmetry A at $Q^2 = 0.48$ GeV^2, where:

$$\begin{aligned} A = \frac{\sigma_R - \sigma_L}{\sigma_R + \sigma_L} &= \left(\frac{-G_F}{\pi\alpha_{em}\sqrt{2}}\right) \frac{1}{\varepsilon\ G_E^2 + \tau\ G_M^2} \\ &\times \left(\varepsilon\ G_E\ G_E^{(Z)} + \tau\ G_M\ G_M^{(Z)} - \frac{1}{2}(1 - 4\sin^2\theta_W)\ \varepsilon'\ G_M\ G_A^{(Z)}\right), \end{aligned} \tag{6}$$

with $\varepsilon = (1 + 2(1+\tau)\tan^2(\theta/2))^{-1}$, $\tau = Q^2/4M^2$, and $\varepsilon' = \sqrt{\tau(1+\tau)(1-\varepsilon^2)}$ (the Q^2 dependence in all form factors is implicit).

Using isospin symmetry, one can relate the electric and magnetic form factors for photon and Z-boson exchange via:

$$G_{E,M}^{(Z)} = \frac{1}{4} G_{E,M}^{(I=1)} - \sin^2\theta_W\ G_{E,M} - \frac{1}{4} G_{E,M}^S, \tag{7}$$

where $G_{E,M}^{(I=1)}$ is the isovector form factor (difference between the proton and neutron). For the $G_{E,M}$ form factors we use the parameterisation from Ref. [44]. The axial form factor for Z-boson exchange is given by $G_A^{(Z)} = -\frac{1}{2}(1 + R_A)G_A + \frac{1}{2}G_A^S$, where R_A is an axial radiative correction, and the axial form factors are known phenomenologically [10].

In Fig.3 we plot the relative difference between the measured asymmetry, A, and that which would be expected for zero strangeness, A_0 — namely, $(A - A_0)/A$. The solid curve corresponds to the light-cone chiral cloud model with a cut-off mass $\Lambda = 1$ GeV for the kaon–hyperon vertex. From the measured HAPPEX asymmetry, the combination $G_E^S + rG_M^S$ was also extracted at an average Q^2 of 0.48 GeV^2, where $r = (\tau/\varepsilon)G_M/G_E \approx 0.4$ for the HAPPEX kinematics. This is shown in Fig.4, compared with the chiral cloud prediction with $\Lambda = 1$ GeV. Both the magnetic (dotted) and electric (dashed) contributions are separately positive, resulting in a small and positive value, consistent with the experiment. Note that exactly the same parameters were used in Figs.3 and 4 as in the fit in Ref. [29] to the SAMPLE data on G_M^S. Therefore the two form factor measurements, as well as the strange axial charge and the strange–antistrange asymmetry, seem to be consistently correlated within the chiral cloud model with soft form factors.

V. CONCLUSION

In this note we have pointed out the existence of corrections to the strange axial charge of the nucleon which arise in light-cone models based on the impulse approximation, or one-body operators, in which Lorentz covariance is not preserved. In the chiral cloud model, where the strangeness content of the nucleon is localised to the kaon–hyperon components of the nucleon wave function, these corrections are an order of magnitude larger than the uncorrected, Lorentz-violating results, and compatible with the sign and magnitude of the empirical g_A^S.

With the same model parameters, namely a soft kaon–hyperon–nucleon form factor (with a kaon probability in the nucleon of $\lesssim 6\%$), one also has good agreement with the strange electromagnetic form factors measured in recent experiments at low Q^2 at MIT-Bates [4] and Jefferson Lab [3]. The results are also compatible with data on the strange–antistrange asymmetry from the CCFR experiments [38].

One should of course mention some of the shortcomings of the simple one-loop meson cloud model treatment, which may qualify some of the quantitative predictions of the model. One of these is the problem of gauge invariance, which in earlier, instant-form approaches has been partially circumvented with the inclusion of contact, or so-called seagull, terms [45]. Unfortunately, these are not unique [46,47], and to date one does not have control over the size of these contributions. Other potential contributions may arise from heavier meson Fock states (such as K^*) or multi-meson configurations. These will be more quantitatively analysed in future work, but our previous experience suggests that their effects are unlikely to be dramatic in a perturbative treatment.

More theoretical work is obviously needed for a deeper understanding of the dynamics of strangeness generation in the nucleon. What seems to be becoming clearer, however, from the accumulating empirical evidence is that the importance of non-perturbative strangeness in the nucleon is likely to be relatively minimal. Future data from Jefferson Lab on the strange electromagnetic form factors, $G_{E,M}^S$, over a range of Q^2 should help to clarify this further.

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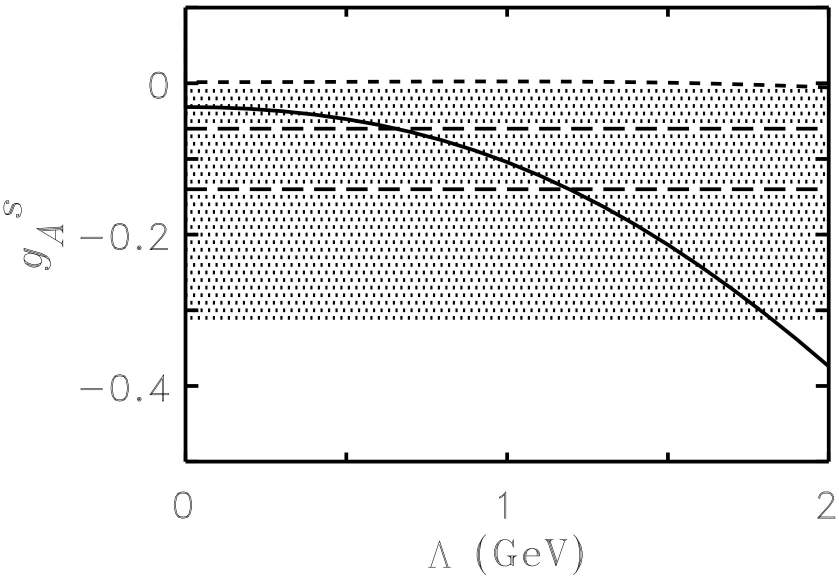

FIG. 1. Strange axial charge of the proton as a function of the hadronic vertex function cut-off mass, Λ. The solid line is the full result from Eq.(4), while the short-dashed is the uncorrected result from Eq.(5). The shaded area represents the range for g_A^S found in νp and $\bar{\nu} p$ elastic scattering [36], and the two long-dashed lines are the limits on Δs from deep-inelastic scattering [37].

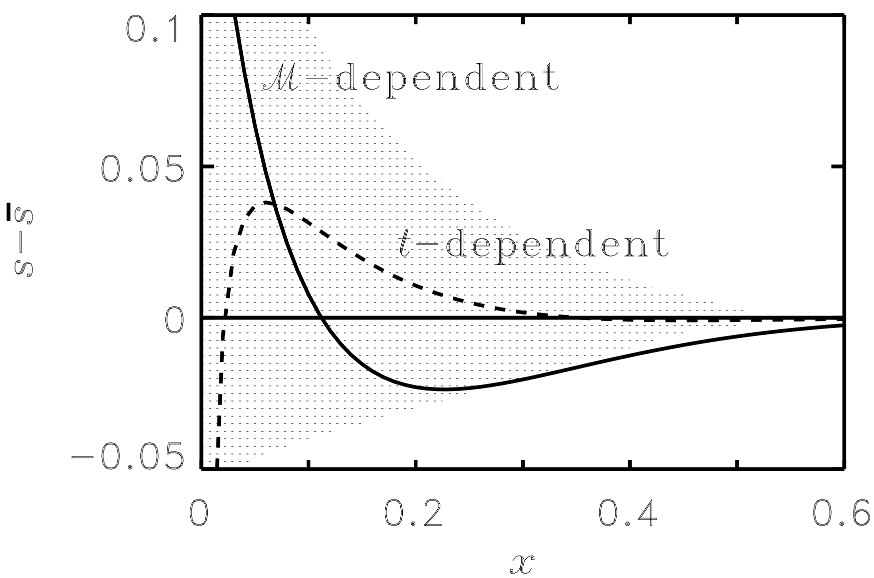

FIG. 2. Strange − antistrange quark difference in the nucleon, with $\mathcal{M}$-dependent (solid) and t-dependent (dashed) monopole form factors, each with a $\Lambda = 1$ GeV momentum cut-off (giving a normalisation of $\langle n \rangle_{KY} \approx 6\%$).

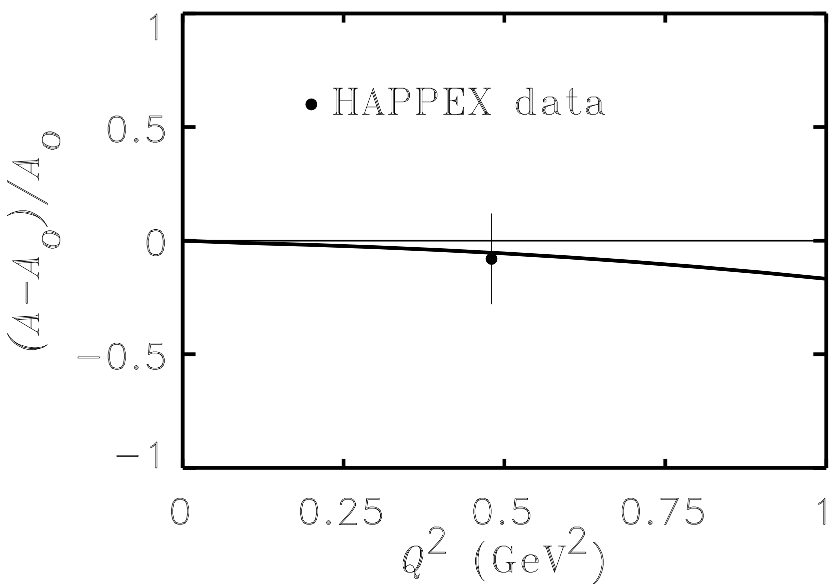

FIG. 3. Relative difference between the measured [3] left-right asymmetry, A, and that expected for zero strangeness, A_0. The solid curve is evaluated with the chiral cloud values for $G^S_{E,M}$ with $\Lambda = 1$ GeV.

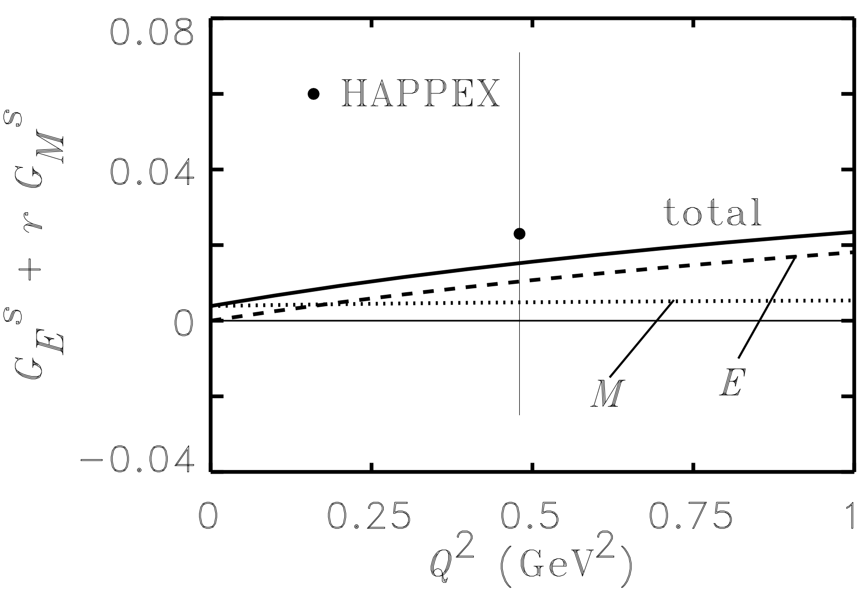

FIG. 4. Strange electric and magnetic form factor combination as extracted from the HAPPEX data [3]. The curves are for the chiral cloud model with a form factor cut-off of $\Lambda = 1$ GeV.